

# Dark Radiation in Anisotropic LARGE Volume Compactifications

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**ABSTRACT:** Dark radiation is a compelling extension to  $\Lambda$ CDM: current experimental results hint at  $\Delta N_{\text{eff}} \gtrsim 0.5$ , which is increased to  $\Delta N_{\text{eff}} \simeq 1$  if the recent BICEP2 results are included. In recent years dark radiation has been considered in the context of string theory models such as the LARGE Volume Scenario of type IIB string theory, forging a link between present-day cosmological observations and models of physics at the Planck scale. In this paper I consider an anisotropic extension of the LARGE Volume Scenario, in which the bulk volume is stabilised by two moduli instead of one. Consequently, the lightest modulus no longer corresponds to the compactification volume but instead to a transverse direction in the bulk geometry. The model also requires sequestering of soft masses, which is achieved by localising the Standard Model on a Euclidean D3 brane wrapping a singularity. I show that the fraction of dark radiation produced in such a model vastly exceeds experimental bounds, ruling out the anisotropic sequestered LARGE Volume Scenario as a model of the early Universe.

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## 1 Introduction

In recent years there has been speculation about the possible existence of an additional relativistic matter component in the energy density of the Universe. This so-called dark radiation is motivated both theoretically and phenomenologically. In UV-complete quantum gravity frameworks such as string theory, the existence of light axion-like particles (ALPs) is commonplace — string compactifications typically produce hundreds of moduli, with associated axions<sup>1</sup> that are massless at the perturbative level due to shift symmetries. Meanwhile, the fact that dark matter is a crucial ingredient in the Standard  $\Lambda$ CDM Cosmological model implores us to ask: if dark matter, then why not dark radiation? The number of relativistic particle species is not protected by any symmetry, therefore there is no reason to assume *a priori* that the present-day radiation content of the Universe must consist of only photons and neutrinos.

Dark radiation is conventionally described in terms of an “excess effective number of neutrino species,”

$$\rho_{\text{DR}} = \frac{7}{8} \left( \frac{4}{11} \right)^{4/3} \rho_{\gamma} \Delta N_{\text{eff}}, \quad (1.1)$$

where  $\Delta N_{\text{eff}} = N_{\text{eff}} - 3.046$ . There are mounting experimental hints for dark radiation. Assuming a tensor-to-scalar ratio  $r = 0$ , a combination of recent CMB observations by Planck [1], high- $l$  data from SPT [2] and ACT [3], WMAP 9-year polarisation data [4], BAO measurements, and the value of  $H_0$  observed by the Hubble Space Telescope [5], suggests

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<sup>1</sup>Hereafter we make liberal use of the term “axion” to refer to axion-like particles in string compactifications.

$N_{\text{eff}} = 3.52^{+0.48}_{-0.45}$  at 95% c.l.. Meanwhile, independent constraints from Big Bang Nucleosynthesis give  $N_{\text{eff}} = 3.50 \pm 0.20$  [6]. The case for dark radiation is further enhanced if one incorporates the recent discovery of primordial B-modes by BICEP2 [7]: using a  $\Lambda\text{CDM}+r$  model with  $r = 0.2^{+0.07}_{-0.05}$ , the authors of [8] find a preference for dark radiation, with  $N_{\text{eff}} = 4.00 \pm 0.41$  [Planck+WP+BICEP2] at 68% c.l.; meanwhile, another preliminary (as-yet unpublished) study [9] points towards values around  $N_{\text{eff}} \sim 3.8 \pm 0.7$  at 95% c.l.. Together, these results provide compelling hints for the possible existence of extra relativistic species.

One of the key motivations for studying dark radiation is that it provides a means of testing models of physics at the Planck scale, such as string theory models. During inflation, the moduli of string compactifications are displaced from their final VEVs, such that when inflation ends they begin to oscillate about their global minimum. Since the moduli behave as non-relativistic matter, they eventually come to dominate the energy density of the Universe. The subsequent reheating of the visible Universe and production of hidden particle species is thus determined by the decay modes of moduli.

In general, moduli have Planck-suppressed decay rates that scale as their mass cubed,

$$\Gamma_{\Phi} \sim \frac{m_{\Phi}^3}{M_{\text{P}}^2}. \quad (1.2)$$

Therefore the lightest modulus is the longest-lived, and since radiation redshifts as  $a^{-4}$  whereas non-relativistic matter evolves as  $a^{-3}$ , any radiation produced by early decays will have redshifted away by the time the lightest modulus decays. Hence reheating is driven solely by the decays of the lightest modulus to the visible sector. Furthermore, this implies that the lightest modulus is also dominantly responsible for dark radiation production.

One phenomenologically appealing string theory model is the LARGE Volume Scenario (LVS) of type IIB string theory [10–12]. In the most basic realisation of this scenario, the overall compactification volume  $\mathcal{V}$  is determined by a single bulk cycle, while additional smaller blow-up cycles can support the visible sector, non-perturbative effects, and additional hidden sectors. The bulk volume is controlled by a Kähler modulus known as the volume modulus, which is stabilised at an exponentially large size due to a combination of  $\alpha'$ -corrections and non-perturbative effects. Consequently, this modulus is hierarchically lighter than all the other moduli, with a mass  $m_{\mathcal{V}} \sim M_{\text{P}}/\mathcal{V}^{3/2}$  (whereas all the other moduli are stabilised around the gravitino mass scale,  $m_{3/2} \sim M_{\text{P}}/\mathcal{V}$ ).

The branching fraction to dark radiation has been studied for this minimal LVS [13, 14], in which the primordial abundance of dark radiation is determined by the decays of the volume modulus to visible- and hidden-sector particles. It turns out that there is one dominant visible-sector decay mode:  $\Phi \rightarrow H_u H_d$  via a Giudice-Masiero term with  $\mathcal{O}(1)$  dimensionless coupling  $Z$  [15].<sup>2</sup> The branching fraction to dark radiation can thus be computed: for the case of a shift symmetry in the Higgs sector, which implies  $Z = 1$  at the string scale [16], one finds a lower bound of  $\Delta N_{\text{eff}} \gtrsim 1.4$ . This is in tension with  $\Delta N_{\text{eff}} \simeq 0.5$  even after loop effects

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<sup>2</sup>This is a dimension-5 operator, so the overall coupling is  $Z/M_{\text{P}}$  times a numerical factor.

are taken into account [17] — this is an example of the “moduli-induced axion problem” [18], which is the statement that string models generically produce too much dark radiation via decays to axion-like particles. However, this tension is relaxed significantly if the BICEP2 results are included in the analysis: a value of  $N_{\text{eff}} = 4.00 \pm 0.41$  [8] is compatible with the minimal LVS, with disagreement at only the  $1\sigma$  level.

It is worthwhile to investigate whether or not extended models can yield a value of  $\Delta N_{\text{eff}}$  that is compatible with observations. One simple extension of LVS is the scenario in which the bulk volume is controlled by two Kähler moduli instead of one [19–22]. Such a setup typically leads to anisotropic modulus stabilisation: one linear combination of these two moduli is the volume modulus, while a transverse flat direction remains unstabilised in the tree-level potential. The crucial feature of anisotropic compactifications most relevant to our purposes is that the volume modulus is no longer the lightest modulus: the post-inflationary decays to visible and hidden radiation are instead controlled by the modulus parametrising the transverse direction. Hence this extension has non-trivial consequences for post-inflationary physics, and one might imagine that the above constraints on dark radiation could thus be avoided. The purpose of this paper is to analyse such a scenario and determine how the branching fraction to dark radiation is modified.

The structure of this paper is as follows. In section 2 we describe and justify an anisotropic compactification scheme, for which we compute the decay modes, and deduce the consequences for  $\Delta N_{\text{eff}}$ , in section 3. In section 4 we conclude.

## 2 Anisotropic compactifications

Here we give an overview of some key features of anisotropic LVS models. First of all, the compactification volume  $\mathcal{V}$  takes the form<sup>3</sup>

$$\mathcal{V} = \alpha \sqrt{\tau_1} \tau_2 - \sum_{i=3}^{h_{1,1}^+} \beta_i \tau_i^{3/2}, \quad (2.1)$$

where  $\tau_1$  and  $\tau_2$  are the Kähler moduli that determine the bulk extra-dimensional volume, and the remaining  $\tau_i$  describe blow-up cycles (“holes”) in the geometry. Such a model will also have  $h_{1,1}^+$  axions  $a_i$ , so we can define complexified Kähler moduli,  $T_i \equiv \tau_i + ia_i$ . In the following section we will neglect all moduli except for  $T_1$  and  $T_2$ , since it turns out that we are focussing on energy scales at which all the other moduli (including complex structure moduli and the axio-dilaton) can be integrated out.

When constructing a realistic model we must bear in mind low-energy phenomenological constraints. In particular, we would like to ensure that soft terms in the visible sector are realised at a scale sufficiently suppressed relative to the masses of all moduli. If this were not the case, requiring TeV-scale superpartners would bring the moduli down to scales  $m_\Phi \lesssim 30 \text{ TeV}$ . Such low moduli masses encounter the Cosmological Moduli Problem (CMP),

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<sup>3</sup>This form arises in fibred compactifications, for example, a K3 or  $T^4$  fibration over a  $\mathbb{P}^1$  base.

in which they dominate the energy density of the Universe at a scale low enough to spoil the successful BBN predictions [26–28].

This problem can be avoided if we consider the case of a Euclidean D3 (ED3) brane wrapping one of the blow-up moduli (say  $\tau_3$ ) and stacks of D7s wrapping  $\tau_1$  and  $\tau_2$ .<sup>4</sup> One obtains a non-perturbative superpotential of the form

$$W = W_0 + Ae^{-aT_3}, \quad (2.2)$$

where  $W_0$  is the tree-level superpotential, which is independent of the Kähler moduli, and the second term is a non-perturbative correction due to instanton effects.

This scenario leads to stabilisation of  $\mathcal{V} \sim \sqrt{\tau_1\tau_2}$  at an exponentially large value, with  $\tau_2 > \tau_1$  [19]. The Standard Model can then be realised on a separate blow-up cycle wrapped by an ED3 brane, leading to a sequestering of the scale at which soft masses appear, as illustrated below.

## 2.1 Mass hierarchy

In this particular anisotropic realisation of LVS, a distinctive hierarchy of mass scales is generated [11, 19]. After diagonalising in terms of mass eigenstates, we find that

$$\begin{aligned} m_{\tau_i} &\sim m_{a_i} \sim \frac{M_P \ln \mathcal{V}}{\mathcal{V}} \quad (i \neq 1, 2), \\ m_{3/2} &\sim m_S \sim m_U \sim \frac{M_P}{\mathcal{V}}, \\ m_\gamma &\sim \frac{M_P}{\mathcal{V}^{3/2}}, \\ m_\Omega &\sim \frac{M_P}{\mathcal{V}^{5/3}}, \\ m_{a_1} &\sim m_{a_2} \simeq 0. \end{aligned} \quad (2.3)$$

Here  $\Omega$  is the direction transverse to the volume modulus. For scenarios in which the Standard Model is located on branes at a singularity [23], soft masses are expected to appear at a scale  $M_{\text{soft}} \sim M_P/\mathcal{V}^2$ ; for TeV-scale superpartners this implies a volume  $\mathcal{V} \sim 5 \times 10^7 \text{ GeV}$ , so  $m_\gamma \sim 3 \times 10^6 \text{ GeV}$  and  $m_\Omega \sim 10^5 \text{ GeV}$ .<sup>5</sup>

## 3 Leading decay modes

We now turn to the computation of the leading decay modes of the lightest modulus  $\Omega$ . In this analysis we focus on the branching fractions to dark radiation and to visible matter, neglecting other possible hidden-sector channels (some of which could also contribute to dark radiation, such as closed-string axions).

<sup>4</sup>This is the “small hierarchy” scenario of [19]. We do not consider the large-hierarchy case since in general soft masses are not sequestered, leading again to the CMP. However, see section 4 for a discussion of an interesting realisation of the large-hierarchy scenario.

<sup>5</sup>If there is no sequestering, the soft masses instead arise at a scale  $m_{3/2} \sim M_P/\mathcal{V}$ , which is incompatible with reheating via the decay of the lightest modulus (which also suffers from the CMP).

### 3.1 Dark radiation

The decay to axions can be computed from the Kähler potential for the bulk Kähler moduli, which can be expressed as

$$K = -\ln(T_1 + \bar{T}_1) - 2\ln(T_2 + \bar{T}_2). \quad (3.1)$$

This is simply the expansion of the usual 4-dimensional  $\mathcal{N} = 1$  supergravity Kähler potential for the Kähler moduli,  $K = -2\ln \mathcal{V}$ , in the anisotropic scenario (up to an irrelevant constant term). From this, we generate un-normalised kinetic terms of the form

$$\mathcal{L} \supset \frac{1}{4\tau_1^2} (\partial_\mu \tau_1 \partial^\mu \tau_1 + \partial_\mu a_1 \partial^\mu a_1) + \frac{1}{2\tau_2^2} (\partial_\mu \tau_2 \partial^\mu \tau_2 + \partial_\mu a_2 \partial^\mu a_2). \quad (3.2)$$

We can canonically normalise the moduli with the reparametrisation

$$\Phi_1 = \frac{1}{\sqrt{2}} \ln \tau_1, \quad \Phi_2 = \ln \tau_2, \quad (3.3)$$

which once expanded out gives

$$\begin{aligned} \mathcal{L} \supset & \frac{1}{2} \partial_\mu \Phi_1 \partial^\mu \Phi_1 + \frac{1}{2} \partial_\mu a_1 \partial^\mu a_2 + \frac{1}{2} \partial_\mu \Phi_2 \partial^\mu \Phi_2 + \frac{1}{2} \partial_\mu a_2 \partial^\mu a_2 \\ & - \sqrt{2} \Phi_1 \partial_\mu a_1 \partial^\mu a_1 - \Phi_2 \partial_\mu a_2 \partial^\mu a_2. \end{aligned} \quad (3.4)$$

The second line of this expression contains the relevant interactions for decays of the  $\Phi$  moduli to axions, which according to (2.3) are massless and therefore constitute dark radiation.

To extract the relevant physics, we must rotate  $\Phi_1$  and  $\Phi_2$  into their mass eigenbasis,

$$\Phi_\mathcal{V} \equiv \sqrt{\frac{2}{3}} \Phi_2 + \sqrt{\frac{1}{3}} \Phi_1, \quad \Phi_\Omega \equiv \sqrt{\frac{1}{3}} \Phi_2 - \sqrt{\frac{2}{3}} \Phi_1, \quad (3.5)$$

where  $\Phi_\mathcal{V}$  is the volume modulus and  $\Phi_\Omega$  is the transverse flat direction. Since  $\Phi_\Omega$  is the lightest modulus its decays will dominate, and hence the relevant interaction for decays into dark radiation is

$$\mathcal{L}_{\Omega \rightarrow aa} = \frac{1}{\sqrt{3} M_{\text{P}}} \Phi_\Omega (2\partial_\mu a_1 \partial^\mu a_1 - \partial_\mu a_2 \partial^\mu a_2). \quad (3.6)$$

This yields a total decay rate to axions of

$$\Gamma_{\Omega \rightarrow aa} = \frac{5}{96\pi} \frac{m_\Omega^3}{M_{\text{P}}^2}, \quad (3.7)$$

which is a factor 5/2 larger than in the minimal LVS.

### 3.2 Visible sector

In the minimal LVS model [13] the leading decay mode is to Higgs bosons via the Giudice-Masiero term [15],

$$\mathcal{L} \supset \frac{1}{\sqrt{6}M_{\text{P}}} \left[ Z H_u H_d \square \Phi + \text{h.c.} \right]. \quad (3.8)$$

Let us see how this is modified in the anisotropic case. We can compute the relevant Lagrangian from a Kähler potential of the form

$$K = -\ln(T_1 + \bar{T}_1) - 2\ln(T_2 + \bar{T}_2) + \left\{ \frac{H_u \bar{H}_u + H_d \bar{H}_d + (Z H_u H_d + \text{h.c.})}{(T_1 + \bar{T}_1)^{1/3} (T_2 + \bar{T}_2)^{2/3}} \right\}. \quad (3.9)$$

This expression contains all the relevant physics: in particular, it incorporates the appropriate scaling of the Kähler matter metric with  $\mathcal{V}^{-2/3}$  [24].

Extracting the leading terms, one finds the interaction Lagrangian

$$\begin{aligned} \mathcal{L} \supset & -\frac{1}{\sqrt{6}M_{\text{P}}} \left( \sqrt{\frac{1}{3}}\Phi_1 + \sqrt{\frac{2}{3}}\Phi_2 \right) [H_u \square \bar{H}_u + \bar{H}_u \square H_u + H_d \square \bar{H}_d + \bar{H}_d \square H_d] \\ & - \frac{1}{\sqrt{6}M_{\text{P}}} (Z H_u H_d + \text{h.c.}) \left( \sqrt{\frac{1}{3}}\square\Phi_1 + \sqrt{\frac{2}{3}}\square\Phi_2 \right). \end{aligned} \quad (3.10)$$

The second line contains the dominant interactions, as the  $\square\Phi$  terms induce a scaling with  $m_{\Phi}^2$ , which from (2.3) is a factor  $\mathcal{V}^{1/2} \sim 10^3$  larger than  $m_H^2$ . The dominant terms have the same structure as (3.8), however, note that the moduli always appear in the combination  $\sqrt{1/3}\Phi_1 + \sqrt{2/3}\Phi_2 \equiv \Phi_{\mathcal{V}}$ . In particular, the Lagrangian is independent of the lightest modulus  $\Phi_{\Omega}$ , so this decay mode is suppressed at tree-level.

Let us consider other possible decay modes of  $\Phi_{\Omega}$ . Chiral matter scalars also interact only with  $\Phi_{\mathcal{V}}$  at tree-level (the relevant Lagrangian has the same form as the first line of (3.10)), while interactions with fermions are chirality-suppressed because they will always contain the Dirac operator  $\bar{\chi}\bar{\sigma}^{\mu}\partial_{\mu}\chi$ , which vanishes on-shell. Furthermore, since the Standard Model is localised on a blow-up cycle, interactions of gauge bosons with the bulk moduli  $\Phi_{\mathcal{V}}$  and  $\Phi_{\Omega}$  will be volume-suppressed, so decay via gauge bosons also takes place only at loop level. Finally, if there are additional vector-like matter states, we expect their tree-level couplings to moduli to be of the same form as (3.10) and hence also independent of  $\Phi_{\Omega}$ . We conclude that all visible-sector decay modes must arise only at loop level, so decays to dark radiation (and possibly other hidden-sector particles) are the dominant processes.

### 3.3 Prediction for the excess effective number of neutrino species

We now provide an estimate for  $\Delta N_{\text{eff}}$  based on these conclusions. Assuming the relevant loop-level decay rates have the form

$$\Gamma_{1\text{-loop}} \sim \left( \frac{\alpha_{\text{SM}}}{4\pi} \right)^2 \frac{m_{\Omega}^3}{M_{\text{P}}^2}, \quad (3.11)$$

where  $\alpha_{\text{SM}}$  represents visible-sector couplings, the ratio of hidden to visible branching ratios is

$$\kappa \equiv \frac{\text{Br}(\text{hidden})}{\text{Br}(\text{visible})} \sim \frac{5\pi}{6} \frac{1}{\alpha_{\text{SM}}^2} \sim 10^2. \quad (3.12)$$

Since  $\Delta N_{\text{eff}} \gtrsim 3\kappa$  [13, 14], this implies  $\Delta N_{\text{eff}} \gtrsim 300$ , which completely rules out the anisotropic scenario with ED3 branes at singularities as a realistic model of dark radiation.

## 4 Conclusions

In this paper we have considered dark radiation in a simple anisotropic extension of the minimal LARGE Volume Scenario. We have focussed on a model in which the visible sector is located on a Euclidean D3 brane at a singularity, which sequesters the soft terms down to order  $M_{\text{P}}/\mathcal{V}^2$ . For a compactification volume  $\mathcal{V} \sim 5 \times 10^7$  the lightest modulus obtains a mass at  $m_{\Omega} \sim 10^5 \text{ GeV}$ , sufficiently heavy to avoid the Cosmological Moduli Problem. We have computed the ratio of branching fractions to hidden- and visible-sector final states and found that this scenario is killed by an excess of dark radiation, with  $\Delta N_{\text{eff}} \gtrsim 300$ . We must therefore turn to other, more elaborate scenarios in order to avoid overproduction of axion-like particles in anisotropic models.

One potentially appealing yet far-fetched possibility is the following. In the scenario considered in [24], poly-instanton corrections from one of the small cycles give rise to a large hierarchy between the bulk moduli  $\tau_1$  and  $\mathcal{V}$ . If a stack of D7-branes wraps  $\tau_1$ , it is possible to achieve a hierarchy of soft masses, with the lightest states separated by a factor  $\mathcal{V}^{-1}$ . One consequence of such a scenario is that the Kähler matter metric would no longer depend only on the bulk volume  $\mathcal{V}$  but also on  $\Omega$ , so a tree-level coupling to Higgs bosons could be restored. Furthermore, since the Standard Model is now located on the very cycle corresponding to the lightest modulus ( $\tau_1$ ), a tree-level coupling to gauge bosons would be generated, further enhancing the branching fraction to the visible sector. Finally, the separation of scales allows for a natural realisation of split supersymmetry, which could explain the lack of observations of supersymmetric particles at the LHC [25].

However, such a scenario suffers from a number of problems. First of all, it is unclear if poly-instantons actually exist in type-IIB string theory. Second, the hierarchical splitting of soft masses in [24] relies on D7s wrapping the fibre cycle  $\tau_1$ . These D7s are likely to generate their own instanton corrections, overwhelming the poly-instantons necessary to generate the anisotropy in the first place. Third, even if such a construction could be realised, it is not clear that it would be stable under loop corrections due to RG running. Such corrections are likely to reduce the hierarchy between states such that an excessively large volume ( $\mathcal{V} \sim 10^{14} \text{ GeV}$ ) is required. This would drive  $m_{\mathcal{V}}$  down to MeV scales, leading once again to the Cosmological Moduli Problem [26–28].

Finally we remark that if the recent discovery of primordial B-modes by BICEP2 [7] holds up, estimates for the effective number of neutrino species could increase  $N_{\text{eff}}$  to values in the region of  $(3.8 \div 4) \pm 0.8$  at 95% confidence level [8, 9]. Such values are significantly



closer to the  $\Delta N_{\text{eff}} \gtrsim 1.4$  prediction of the minimal LVS, so ultimately the minimal scenario may be saved after all.

*NOTE: This paper is submitted simultaneously to the related work [29].*

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